Gauge Mediation Simplified

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Gauge mediation of supersymmetry breaking is drastically simplified using generic superpotentials without $U(1)_R$ symmetry by allowing metastable vacua.

Breaking supersymmetry has been a non-trivial task. A general argument by Nelson and Seiberg is that it requires a theory with a continuous exact $U(1)_R$ symmetry if we assume that the superpotential is generic [1]. In addition, an argument based on the Witten index [2] said that the theory must be chiral. This is because one can continuously deform a vector-like theory by mass terms to a pure Yang-Mills theory, which is known to have a finite Witten index (dual Coxeter number) and hence supersymmetric vacua. Chirality and $U(1)_R$ invariance strongly limit the choice of possible theories that break supersymmetry. Therefore explicit models of supersymmetry breaking appear rather special and hence do not seem likely to come out from a more fundamental theory such as string theory. This problem is exacerbated by the fact that the supersymmetry breaking sector should couple to the standard model multiplets to induce soft supersymmetry breaking parameters in a flavor-independent fashion.

Later, vector-like models were found [3]. They evade the Witten index argument because the mass terms can always be absorbed by shifting singlet fields in the theory. The required superpotential, however, is not generic unless one imposes an exact $U(1)_R$ symmetry.

The requirement of an exact $U(1)_R$ symmetry is unfortunate, because exact global symmetries are not expected to exist in quantum theory of gravity such as the field-theory limit of string theory. In addition, embedding a model of supersymmetry breaking into supergravity requires explicit breaking of $U(1)_R$ to allow for a constant term in the superpotential needed for canceling the cosmological constant. Once $U(1)_R$ is not an exact symmetry, it is not clear how one can justify the form of the superpotential required for supersymmetry breaking.

In this letter, we advocate to discard $U(1)_R$ symmetry altogether from the theory, and allow for completely generic superpotentials. According to the Nelson–Seiberg argument, such a theory would not break supersymmetry. Yet, it may have a *local* supersymmetry breaking minimum. Supersymmetry is broken if the low-energy limit of the supersymmetry breaking sector has an *accidental* $U(1)_R$ symmetry, which nonetheless is broken by its coupling to messengers. Indeed, we show a very simple class of models of this type. The models do not have a fundamental singlet field, eliminating aesthetic and various fine-tuning problems in cosmology and preserving the hierarchy. The gauginos and scalars in the supersymmetric standard model sector obtain flavor universal masses by standard model gauge interactions through loops of the messengers. Given the absence of $U(1)_R$, there is no problem in generating gaugino masses, and no dangerous R-axion arises.

An explicit model that realizes our general philosophy is a supersymmetric $SU(N_c)$ QCD with massive vector-like quarks Q^i and $\bar{Q}^i$ ($i = 1, \cdots, N_f$). In addition, we introduce massive messengers f and $\bar{f}$ and write the most general superpotential consistent with the gauge symmetry. This is the entire model. The important terms in the superpotential are given by

$$W_{\rm tree} = m_{ij}\bar{Q}^i Q^j + \frac{\lambda_{ij}}{M_{\rm Pl}}\bar{Q}^i Q^j \bar{f} f + M\bar{f}f, \qquad (1)$$

where λ_{ij} are coupling constants [12]. (The effects of other terms will be discussed later.) For concreteness, we take the messengers $f, \bar{f}$ to be in $\mathbf{5}+\mathbf{5}^*$ representations of $SU(5)$ in which the standard model gauge group is embedded.

Intriligator, Seiberg, and Shih (ISS) pointed out that supersymmetric $SU(N_c)$ QCD in the free magnetic phase ($N_c+1 \leq N_f < \frac{3}{2}N_c$) breaks supersymmetry on a metastable local minimum if the quark masses m_{ij} are much smaller than the dynamical scale Λ [4]. Note that in the ISS model a $U(1)_R$ symmetry is broken only down to Z_{2N_c} which prevents the gaugino masses. In the present model, however, the coupling to the messengers breaks it down to Z_2, so that the model does not have any R symmetry beyond R-parity.

For the sake of concreteness, we discuss the case without the magnetic gauge group $N_f = N_c+1$ below, although any $N_c+1 \leq N_f < \frac{3}{2}N_c$ works equally well. At energies below the dynamical scale, the non-perturbative low-energy effective superpotential is described as [5]

$$W_{\rm dyn} = \frac{1}{\Lambda^{2N_f-3}}\left(\bar{B}_i M^{ij} B_j - {\rm det} M^{ij}\right), \qquad (2)$$

where $M^{ij} = \bar{Q}^i Q^j$, $B_i = \epsilon_{ii_1\cdots i_{N_c}} Q^{i_1}\cdots Q^{i_{N_c}}/N_c!$ and $\bar{B}_i = \epsilon_{ii_1\cdots i_{N_c}}\bar{Q}^{i_1}\cdots\bar{Q}^{i_{N_c}}/N_c!$ are meson, baryon and antibaryon chiral superfields, respectively. In the following, we adopt the basis in which the quark mass matrix is diagonal, $m_{ij} = -m_i\delta_{ij}$, with m_i real and positive. We also assume that they are ordered as $m_1 > m_2 > \cdots > m_{N_f} > 0$ without loss of generality. Here, we have taken all masses different to avoid (potentially) unwanted Nambu–Goldstone bosons.

In terms of fields with canonical dimensions $S^{ij} = M^{ij}/\Lambda$, $b_i = B_i/\Lambda^{N_f-2}$ and $\bar{b}_i = \bar{B}_i/\Lambda^{N_f-2}$, the dynamical superpotential of Eq. (2) together with the quark mass terms (the first term of Eq. (1)) can be written as [13]

$$W_{\rm ISS} = \bar{b}_i S^{ij} b_j - \frac{{\rm det} S^{ij}}{\Lambda^{N_f-3}} - m_i \Lambda S^{ii}. \tag{3}$$

For $N_f > 3$, the superpotential term ${\rm det}S^{ij}$ is irrelevant and can be ignored to discuss physics around the origin $S^{ij} = 0$ [14]. The superpotential of Eq. (3) then leads to a local minimum at

$$b = \bar{b} = \begin{pmatrix} \sqrt{m_1 \Lambda} \\ 0 \\ \vdots \\ 0 \end{pmatrix}, \qquad S^{ij} = 0, \tag{4}$$

where supersymmetry is broken because $F_{S^{ij}} = -(\partial_{S^{ij}} W)^* = m_i \delta_{ij} \Lambda \neq 0$ for $i, j \neq 1$. Even though S^{ij} ($i, j \neq 1$) are classically flat directions, they are lifted by the one-loop Coleman–Weinberg potential. As a result, the origin $S^{ij} = 0$ is a local minimum, with curvature $m_{S^{ij}}^2 \sim m\Lambda/16\pi^2$ for all $m_i \sim m$. It is long-lived as long as $m_i \ll \Lambda$, where the weakly-coupled analysis of the low-energy theory is valid.

The existence of a supersymmetry breaking minimum of Eq. (4) can be viewed as a result of an accidental (and approximate) $U(1)_R$ symmetry possessed by the superpotential of Eq. (3) with the R-charge assignments $R(S^{ij}) = 2$, $R(b_i) = R(\bar{b}_i) = 0$, in the limit of neglecting the irrelevant term of ${\rm det}S^{ij}/\Lambda^{N_f-3}$. In fact, this accidental $U(1)_R$ symmetry is also a reason for the origin $S^{ij} = 0$ being the minimum of the effective potential as a symmetry enhanced point. This picture is corrected by the coupling of Q^i and $\bar{Q}^i$ to the messengers and by higher dimension terms in the superpotential omitted in Eq. (1), which introduce $U(1)_R$ violating effects to the supersymmetry breaking sector. These effects, however, can be easily suppressed as we will see later, and the basic picture described above can be a good approximation of the dynamics.

At the supersymmetry breaking minimum of Eq. (4) (with S^{ij} slightly shifted due to $U(1)_R$ violating effects), the messenger fields have both supersymmetric and holomorphic supersymmetry breaking masses:

$$M_{\rm mess} = M + \frac{\lambda_{ij}\Lambda}{M_{\rm Pl}} \langle S^{ij} \rangle \simeq M, \tag{5}$$

and

$$F_{\rm mess} = \frac{\lambda_{ij}\Lambda}{M_{\rm Pl}} F_{S^{ij}} = \frac{\bar{m}\Lambda^2}{M_{\rm Pl}}, \tag{6}$$

where

$$\bar{m} \equiv \sum_{i \neq 1} \lambda_{ii} m_i. \tag{7}$$

The usual loop diagrams of the messenger fields then induce gauge-mediated scalar and gaugino masses in the supersymmetric standard model sector, of the magnitude [6, 7]

$$m_{\rm SUSY} \simeq \frac{g^2}{16\pi^2} \frac{\bar{m}\Lambda^2}{M M_{\rm Pl}}, \tag{8}$$

where g represents generic standard model gauge coupling constants.

Several conditions for the parameters need to be met for the model to be phenomenologically successful. Even though not necessary, we regard all the quark masses (and the couplings λ_{ij}) to be comparable, $m_i \sim m$ ($\lambda_{ij} \sim \lambda$), in the numerical estimates below.

First, we would like $m_{\rm SUSY}$ to stabilize the electroweak scale, and hence $m_{\rm SUSY} = O(100~{\rm GeV} \sim 1~{\rm TeV})$. This corresponds to

$$\frac{\bar{m}\Lambda^2}{M M_{\rm Pl}} \approx 100~{\rm TeV}. \tag{9}$$

On the other hand, we would like the gauge-mediated contribution to the scalar masses dominate over the gravity-mediated piece to avoid excessive flavor-changing processes, leading to $m_{3/2} \approx m\Lambda/M_{\rm Pl} \lesssim 10^{-2} m_{\rm SUSY}$. Therefore,

$$m M \lesssim 10^{-4} \bar{m}\Lambda. \tag{10}$$

We also need the messengers to be non-tachyonic,

$$M^2 > \frac{\bar{m}\Lambda^2}{M_{\rm Pl}}. \tag{11}$$

In addition, the analysis of supersymmetry breaking is valid only if m is sufficiently smaller than Λ:

$$m \lesssim 0.1\Lambda. \tag{12}$$

We now discuss the effects of $U(1)_R$ violation. These effects cause shifts of S^{ij} from the origin, which must be smaller than $\approx 4\pi\sqrt{m\Lambda}$ for the ISS analysis to be valid, and than $\approx M M_{\rm Pl}/\lambda\Lambda$ to avoid tachyonic messengers. One origin of $U(1)_R$ violation comes from higher dimension terms in the superpotential, omitted in Eq. (1). The dominant effect comes from

$$\Delta W = \frac{\lambda_{ijkl}}{M_{\rm Pl}} \bar{Q}^i Q^j \bar{Q}^k Q^l = \frac{\lambda_{ijkl}\Lambda^2}{M_{\rm Pl}} S^{ij} S^{kl}. \tag{13}$$

These terms may destabilize the minimum, since they lead to linear terms of S^{ij} in the potential through $F_{S^{ij}} = m_i \delta_{ij} \Lambda$ [8]. The squared masses of S^{ij} from the one-loop effective potential are $m_{S^{ij}}^2 \sim m\Lambda/16\pi^2$, while the linear terms are $\sim (\lambda_{ijkk} m_k \Lambda^3/M_{\rm Pl}) S^{ij}$. Therefore, the shifts of the fields are $\Delta S^{ij} \sim 16\pi^2 \lambda_{ijkk}\Lambda^2/M_{\rm Pl}$. Requiring this to be sufficiently small, we obtain the condition

$$\frac{\lambda_{ijkk}\Lambda^2}{M_{\rm Pl}} \lesssim \min\left\{ 0.1 (m\Lambda)^{1/2}, 10^{-2} \frac{M M_{\rm Pl}}{\lambda\Lambda} \right\}. \tag{14}$$

Similar conditions can be worked out for even higher order terms, but they are rather mild.

Another source of $U(1)_R$ violation comes from the coupling of Q^i and $\bar{Q}^i$ to the messengers, which shifts the minimum of S^{ij} at the loop level. The effect of the messengers on the S^{ij} effective potential can be calculated by computing the one-loop Coleman–Weinberg potential arising from the last two terms of Eq. (1):

$$W_{\rm mess} = \frac{\lambda_{ij}\Lambda}{M_{\rm Pl}} S^{ij} \bar{f} f + M \bar{f} f. \tag{15}$$

The resulting effective potential takes the following generic form

$$\Delta V \approx \frac{\bar{m}^2 \Lambda^4}{16\pi^2 M_{\rm Pl}^2} \mathcal{F}\left(\frac{\lambda_{ij}\Lambda S^{ij}}{M M_{\rm Pl}}\right), \tag{16}$$

where $\mathcal{F}(x)$ is a real polynomial function with the coefficients of $O(1)$ up to symmetry factors. The resulting shifts of S^{ij} are of order $\lambda^3 m \Lambda^4 / M M_{\rm Pl}^3$, which are sufficiently small if

$$M \gtrsim \frac{\lambda^2 m^{1/2} \Lambda^{5/2}}{M_{\rm Pl}^2}. \tag{17}$$

Note that the coupling to the messengers in Eq. (15) does not generate a new supersymmetric minimum. However, turning on the expectation values for the messengers may allow for lowering the vacuum energy, depending on the combinations of m_{ij} and $\lambda_{ij}\bar{f}f$. Even if this is the case, the tunneling to a lower minimum at $\bar{f}f \approx m M_{\rm Pl}/\lambda$ can easily be made suppressed to the level consistent with the longevity of our universe, if $M M_{\rm Pl}/\lambda \gtrsim m^{1/2}\Lambda^{3/2}$.

It is now easy to see that there is a wide range of parameters that satisfy the conditions Eqs. (9, 10, 11, 12, 14, 17). For instance, if we take $\lambda_{ij} \sim \lambda_{ijkl} \sim 1$, $\Lambda \sim 10^{11}$ GeV, $m \sim \bar{m} \sim 10^8$ GeV and $M \sim 10^7$ GeV, then all the requirements are easily satisfied. Note that the conditions of Eqs. (14, 17) are generically rather weak, unless Λ is close to $M_{\rm Pl}$. This is because the relevant interactions in Eqs. (13, 15) arise from higher dimension operators suppressed by $M_{\rm Pl}$.

Finally, we discuss if there are any unwanted light fields in the model. The fermionic fields in S^{ij} $(i, j \neq 1)$ are massless in the ISS model, but they acquire masses here due to the generic terms in Eq. (13) [15]. They can decay to standard model particles through their coupling to the messengers and hence harmless. There is a Nambu–Goldstone boson (NGB) of a spontaneously broken $U(1)_B$ symmetry, $b^1 - \bar{b}^1$, and its fermionic partner. Exactly massless NGB and fermion would be a radiation component of the universe. Their abundance is diluted by an order of magnitude due to the QCD phase transition and is in general consistent with the constraint from the big-bang nucleosynthesis, $\Delta N_\nu \lesssim 1.5$ [9]. Alternatively, they can be made massive by gauging $U(1)_B$, or avoided entirely by employing an $SO(N_c)$ or $Sp(N_c)$ gauge group for supersymmetry breaking, instead of $SU(N_c)$. The gravitino is the lightest supersymmetric particle and hence stable if R-parity is unbroken. It places an upper limit on the reheating temperature [10], which is acceptable *e.g.*, in leptogenesis models by non-thermal production of right-handed scalar neutrinos [11].

In summary, we advocated gauge mediation models of supersymmetry breaking with generic superpotentials without $U(1)_R$ symmetry. Using metastable minima, we find a class of phenomenologically successful models without any elementary gauge singlet fields. We find the simplicity and generality of the models quite remarkable.

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[12] Here, we took the scale of higher dimension operators to be the reduced Planck scale $M_{\rm Pl}$ just for the sake of presentation, but of course it can be some other scales as

well.
[13] The fields S^{ij}, b_i and $\bar{b}_i$ are in general not canonically normalized by incalculable $O(1)$ wavefunction renormalization factors, which are not important to our discussions and hence disregarded in the rest of the letter.
[14] This term, however, is important to see that there are global supersymmetric minima at nonzero S^{ij} as suggested by the general arguments.
[15] Of course, one of these fields remains massless as the Goldstino which is eaten by the gravitino.